

Rare top quark decay $t \rightarrow u_1 \bar{u}_2 u_2$ in the standard model

A. Cordero-Cid, J. M. Hernández, G. Tavares-Velasco, and J. J. Toscano
*Facultad de Ciencias Físico Matemáticas, Benemérita Universidad
 Autónoma de Puebla, Apartado Postal 1152, Puebla, Pue., México*
 (Dated: February 2, 2008)

The one-loop induced top quark decay $t \rightarrow u_1 \bar{u}_2 u_2$ ($u_i = u, c$) is calculated in the context of the standard model. The dominant contribution to this top quark decay arises from the Feynman diagrams induced by the off-shell $tu_1 g^*$ vertex, whereas the box diagrams are negligibly small. In contrast with the on-shell $tu_1 g$ vertex, which only gives rise to a pure dipolar effect, the off-shell $tu_1 g^*$ coupling also involves a monopolar term, which gives a larger contribution than the dipolar one. It is found that the branching ratio for the three-body decay $t \rightarrow u_1 \bar{u}_2 u_2$ is about of the same order of magnitude of the two-body decay $t \rightarrow u_1 g$, which stems from the fact that the three-body decay is dominated by the monopolar term.

PACS numbers: 14.65.Ha, 12.15.Ji, 12.15.Lk

I. INTRODUCTION

The top quark detection at the Fermilab Tevatron [1] greatly boosted the interest in top quark physics. The large mass of this quark suggests that it could be very sensitive to new physics effects, which may manifest themselves through anomalous rates for its production and decay modes. Although some properties of the top quark have already been examined at the Tevatron [2], a further scrutiny is expected at the CERN large hadron collider (LHC). This machine will operate as a veritable top quark factory, producing about eight millions of $t\bar{t}$ events per year in its first stage, and hopefully up to about eighteen millions in subsequent years [3]. Yet in the first stage of the LHC, many rare processes involving the top quark are expected to be accessible. It is thus worth investigating all of the top quark decays within the standard model (SM) in order to find out any scenario that may be highly sensitive to new physics effects.

In the SM, the main decay channel of the top quark is $t \rightarrow bW$. Although the nondiagonal $t \rightarrow dW$ and $t \rightarrow sW$ modes are more suppressed, they still have sizable branching ratios. For instance, $Br(t \rightarrow sW)$ is of the order of 10^{-3} . As far as rare decays are concerned, the three-body tree-level induced modes $t \rightarrow d_i WZ$ and $t \rightarrow u_1 WW$, with $d_i = b, s, d$ and $u_1 = u, c$, are strongly dependent on the precise value of the top quark mass. It has been shown that the $t \rightarrow u_1 WW$ decays are severely GIM-suppressed [4], but $t \rightarrow bWZ$ can have a branching ratio of the order of 10^{-5} for a top quark mass larger than 187 GeV [5]. This decay mode has been suggested as a probe for the top quark mass because it is almost in the threshold region [6]. At the one-loop level, there arise the flavor changing neutral current (FCNC) decays $t \rightarrow u_1 V$ ($V = g, \gamma, Z$) and $t \rightarrow u_1 H$, which are considerably GIM-suppressed, with branching ratios ranging from 10^{-10} to 10^{-13} [7, 8, 9]. Motivated by the fact that any process that is forbidden or strongly suppressed within the SM constitutes a natural laboratory to search for any new physics effects, FCNC top quark decays have been the subject of considerable interest in the literature [10, 11, 12, 13, 14]. It turns out that they may have large branching ratios, much larger than the SM ones, within some extended theories such as the two-Higgs doublet model (THDM) [10], supersymmetry (SUSY) models with nonuniversal soft breaking [11], SUSY models with broken R -parity [12], and even more exotic scenarios [13]. Similar results for the decays $t \rightarrow u_1 V$ and $t \rightarrow u_1 H$ were obtained within the context of effective theories [14].

In this work, we present a calculation of the $t \rightarrow u_1 \bar{u}_2 u_2$ decay (u_2 stands for the u or c quark), which arises at the one-loop level in the SM. Although the study of rare top quark transitions has attracted considerable attention, to our knowledge the rare decay $t \rightarrow u_1 \bar{u}_2 u_2$ has never been analyzed before. The rest of the paper is organized as follows. Section II is devoted to the analytical calculation of the decay $t \rightarrow u_1 \bar{u}_2 u_2$. The numerical results and discussion are presented in Sec. III along with the final remarks.

II. THE DECAY $t \rightarrow u_1 \bar{u}_2 u_2$

Decays of the type $t \rightarrow u_1 \bar{u}_2 u_2$ proceed through the reducible diagrams shown in Fig. 1(i), which are mediated by the Z , H , γ or g bosons. While those Feynman diagrams mediated by the Z and H bosons are enhanced due to the fact that the intermediary boson is on resonance (provided that $m_H \leq m_t$), those diagrams mediated by the photon (gluon) are enhanced by the effect of the photon (gluon) pole. There are also contributions from box diagrams carrying W gauge bosons and down quarks [see Fig. 1(ii)]. Since each type of diagram renders a finite amplitude

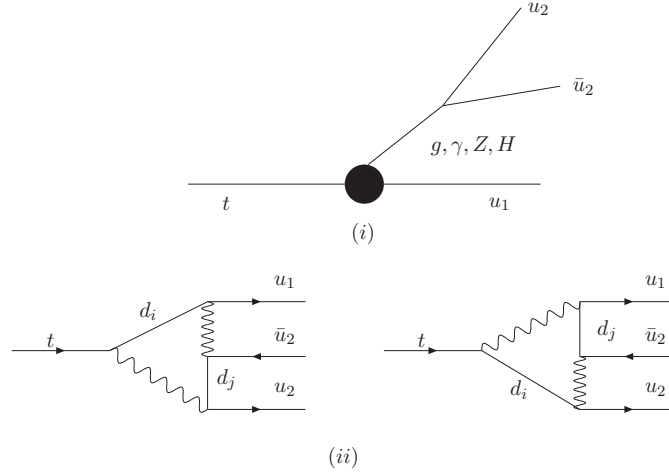


FIG. 1: Feynman diagrams contributing to the $t \rightarrow u_1 u_2 \bar{u}_2$ decay. The bubble stands for all the contributions of the type shown in Fig. 2.

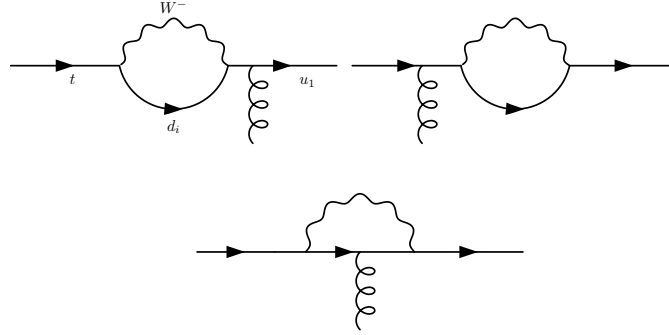


FIG. 2: Feynman diagrams contributing to the tu_1g vertex in the unitary gauge. In the R_ξ gauge there are an extra set of diagrams in which the W boson is replaced by its associate pseudogoldstone boson.

by its own, the different contributions can be considered as independent. It can be shown that the $t \rightarrow u_1 \bar{u}_2 u_2$ decay is essentially determined by those graphs involving a virtual gluon, i.e., those reducible diagrams involving the one-loop vertex tu_1g^* and the tree-level vertex $g^*\bar{u}_2u_2$. The role played by this contribution is evident from the fact that it constitutes an electroweak-QCD mixed effect. This is to be contrasted with those reducible graphs mediated by the γ , Z and H bosons, as well as the box diagrams, which are entirely determined by electroweak couplings. As a consequence, the pure electroweak contributions become suppressed by a factor of α/α_s as compared with the electroweak-QCD mixed ones.

Once the most relevant properties of the $t \rightarrow u_1 \bar{u}_2 u_2$ decay were described, we would like to emphasize a noteworthy feature of this process. It is closely related to the fact that the $t \rightarrow u_1 \bar{u}_2 u_2$ decay is mediated by a virtual massless vector boson, i.e., the gluon or the photon. Without losing generality, it is enough to discuss the gluon contribution as it is the dominant one. Naively, one would expect that the rate for the two-body decay $t \rightarrow u_1 g$ is larger than that for the three-body decay $t \rightarrow u_1 \bar{u}_2 u_2$, which in fact is not true. This stems from the fact that while the on-shell tu_1g vertex is characterized by a dipole structure (the tu_1 pair couples to the gluon through the gauge tensor $G_{\mu\nu}^a$), the corresponding off-shell tu_1g^* vertex involves also a monopole structure (the tu_1 pair interacts directly with the A_μ^a gauge field). Therefore, while the $t \rightarrow u_1 g$ transition is entirely determined by the dipole structure, both the dipole and the monopole structures contribute to the $t \rightarrow u_1 \bar{u}_2 u_2$ process. It turns out that the contribution from the monopolar term can be considerably larger than that arising from the dipolar one. We have found that this is indeed the case for the rare decay $t \rightarrow u_1 \bar{u}_2 u_2$. It means that while the $t \rightarrow u_1 g$ decay is determined by the dipolar term, the $t \rightarrow u_1 \bar{u}_2 u_2$ mode is governed by the monopolar one. Moreover, the three-body decay is unsuppressed because it includes the QCD vertex $g^*\bar{u}_2u_2$, which is much less suppressed than the electroweak vertices $\gamma^*\bar{u}_2u_2$ and $Z^*\bar{u}_2u_2$. The above properties nicely conspire to enhance the $t \rightarrow u_1 \bar{u}_2 u_2$ decay rate by about one order of magnitude with respect to that of the $t \rightarrow u_1 g$ transition.

We now turn to analyze the general structure of the tu_1g vertex, which is generated at the one-loop level via the

Feynman diagrams shown in Fig. 2. The most general form of this vertex involves up to ten form factors, which are associated with the Lorentz structures $\gamma_\mu P_{L,R}$, $(p - p_1)_\mu P_{L,R}$, $q_\mu P_{L,R}$, $\sigma_{\mu\nu}(p - p_1)^\nu P_{L,R}$, and $\sigma_{\mu\nu}q^\nu P_{L,R}$, with $P_{L,R} = (1 \mp \gamma_5)/2$. There are however a few independent form factors. After imposing the on-shell conditions on the fermionic fields, the Gordon identities allows one to eliminate the four form factors associated with $(p - p_1)_\mu P_{L,R}$ and $\sigma_{\mu\nu}(p - p_1)^\nu P_{L,R}$. In addition, the $q_\mu \epsilon^\mu = 0$ condition, valid for a real gluon, can be safely used as the off-shell gluon couples to a pair of approximately massless quarks. Thus, the only non negligible contributions to the $t \rightarrow u_1 \bar{u}_2 u_2$ decay are those of the monopolar ($\gamma_\mu P_{L,R}$) and dipolar ($\sigma_{\mu\nu}q^\nu P_{L,R}$) terms. Furthermore, no contributions proportional to P_R can arise in the $m_{u_1} = 0$ limit. Consequently, the vertex function associated with the tu_1g coupling can be written as

$$\Gamma_\mu^a = \frac{\lambda^a}{2} \Gamma_\mu = \left(\frac{\lambda^a}{2} \right) i g_s \left[F_1(q^2) \gamma_\mu P_L + \frac{i}{m_t} F_2(q^2) P_L \sigma_{\mu\nu} q^\nu \right], \quad (1)$$

where g_s is the coupling constant and $\lambda^a/2$ stand for the generators associated with the color group. It is worth mentioning that the monopolar contribution F_1 vanishes in the on-shell limit as a consequence of the Ward identity $q^\mu \Gamma_\mu = 0$. This means that a FCNC vertex involving an on-shell gluon can only arise through a dipolar term, such as occurs in electrodynamics. This is not true for an off-shell gluon: in such a case the monopolar term yields the dominant contribution. This behavior will be explicitly shown below. We will first calculate the form factors in an R_ξ gauge and, in order to assure that our result is gauge-independent, we will also calculate such form factors via the unitary gauge. In the R_ξ gauge the calculation leads to the following amplitude

$$\Gamma_\mu = \frac{g_s g^2}{2} \sum_i V_{u_1 i} V_{it}^\dagger \int \frac{d^D k}{(2\pi)^D} \sum_{a=1}^6 \frac{T_\mu^a}{\Delta_a}, \quad (2)$$

where

$$T_\mu^1 = \gamma_\beta P_L (\not{k} + \not{p}_2 + m_i) \gamma_\mu (\not{k} + \not{p}_1 + m_i) \gamma_\alpha P_L P^{\alpha\beta}, \quad (3)$$

$$T_\mu^2 = - \left(\frac{1}{m_t^2} \right) \gamma_\beta P_L (\not{k} + \not{p}_2 + m_i) \gamma_\alpha P_L (\not{p}_2 + m_t) \gamma_\mu P^{\alpha\beta}, \quad (4)$$

$$T_\mu^3 = \left(\frac{1}{m_t^2} \right) \gamma_\mu \not{p}_1 \gamma_\beta P_L (\not{k} + \not{p}_1 + m_i) \gamma_\alpha P_L P^{\alpha\beta}, \quad (5)$$

$$T_\mu^4 = \left(\frac{m_i}{m_W^2} \right) P_R (\not{k} + \not{p}_2 + m_i) \gamma_\mu (\not{k} + \not{p}_1 + m_i) (m_t P_R - m_i P_L), \quad (6)$$

$$T_\mu^5 = - \left(\frac{1}{m_t^2} \right) \left(\frac{m_i}{m_W^2} \right) P_R (\not{k} + \not{p}_2 + m_i) (m_t P_R - m_i P_L) (\not{p}_2 + m_t) \gamma_\mu, \quad (7)$$

$$T_\mu^6 = \left(\frac{1}{m_t^2} \right) \left(\frac{m_i}{m_W^2} \right) \gamma_\mu \not{p}_1 (\not{k} + \not{p}_1 + m_i) (m_t P_R - m_i P_L), \quad (8)$$

$$\Delta_1 = [k^2 - m_W^2][(k + p_1)^2 - m_i^2][(k + p_2)^2 - m_i^2], \quad (9)$$

$$\Delta_2 = [k^2 - m_W^2][(k + p_2)^2 - m_i^2], \quad (10)$$

$$\Delta_3 = [k^2 - m_W^2][(k + p_1)^2 - m_i^2], \quad (11)$$

$$\Delta_4 = [k^2 - \xi m_W^2][(k + p_1)^2 - m_i^2][(k + p_2)^2 - m_i^2], \quad (12)$$

$$\Delta_5 = [k^2 - \xi m_W^2][(k + p_2)^2 - m_i^2], \quad (13)$$

$$\Delta_6 = [k^2 - \xi m_W^2][(k + p_1)^2 - m_i^2], \quad (14)$$

and

$$P^{\alpha\beta} = g^{\alpha\beta} - \frac{(1 - \xi) k^\alpha k^\beta}{k^2 - \xi m_W^2}. \quad (15)$$

In the previous expressions, the $m_{u_1} = 0$ approximation was used. In addition, m_i denotes the mass of the internal down quark and ξ is the gauge parameter. For simplicity the calculation was performed in the t'Hooft-Feynman gauge ($\xi = 1$). As a crosscheck, we also have performed the calculation via the unitary gauge ($\xi = \infty$), in which there are only contributions from the first three terms in Eq. (2), with $P^{\alpha\beta}$ replaced by $P^{\alpha\beta} = g^{\alpha\beta} - k^\alpha k^\beta / m_W^2$. The results obtained by these two calculation schemes do coincide, which guarantees that the form factors associated with the monopolar and dipolar structures of the $tu_1 g^*$ vertex are gauge-independent. Introducing the definition $F_i = (\alpha^2/8\pi)A_i$ we can write

$$A_1 = \frac{x}{2x_W(1-x)} \sum_{i=d,s,b} V_{u_1 i} V_{ti}^\dagger \left(f_0^1 + \sum_{a=1}^3 f_a^1 B_0(a) + 2m_t^2 g_1 C_0(x, x_i) \right), \quad (16)$$

$$A_2 = \frac{1}{2x_W(1-x)} \sum_{i=d,s,b} V_{u_1 i} V_{ti}^\dagger \left(f_0^2 + \sum_{a=1}^3 f_a^2 B_0(a) + 2m_t^2 g_2 C_0(x, x_i) \right), \quad (17)$$

where $x = q^2/m_t^2$, $x_W = m_W^2/m_t^2$, and $x_i = m_i^2/m_t^2$. The f_a^b functions depend on x_i and read

$$f_0^1 = x_i, \quad (18)$$

$$\begin{aligned} f_1^1 = & \frac{1}{1-x} \left(\left(2(1-x) + (x-4)(x_i + x_W) \right) x_i \right. \\ & \left. - 2 \left(2(1-x) + (x-4)x_W \right) x_W \right), \end{aligned} \quad (19)$$

$$\begin{aligned} f_2^1 = & \frac{1}{(1-x)^2} \left((x+2)(1-x-2(x_i + x_W)) x_i \right. \\ & \left. + 2 \left(2(x+2)x_W - x(1-x) - 2(1-x^2) \right) x_W \right) \end{aligned} \quad (20)$$

$$\begin{aligned} f_3^1 = & \frac{1}{(1-x)^2} \left(\left(-x^2 + 5x - 4 + (x^2 - 3x + 8)(x_i + x_W) \right) x_i \right. \\ & \left. - 2 \left(x^2 + 3x - 4 + (x^2 - 3x + 8)x_W \right) x_W \right), \end{aligned} \quad (21)$$

$$\begin{aligned} g_1 = & \frac{1}{(1-x)^2} \left(2 \left(x(1-x)^2 + (x+2)x_W^2 + 2(x^2 - 1)x_W \right) x_W \right. \\ & \left. + \left((1-x)^2(x_i - 1) + (x+2)(x_i^2 - 3x_W^2) - 2((x+1)^2 - 4)x_W \right) x_i \right), \end{aligned} \quad (22)$$

and

$$f_0^2 = x_i, \quad (23)$$

$$\begin{aligned} f_1^2 = & \frac{1}{1-x} \left(\left(2(1-x) - (x+2)(x_i + x_W) \right) x_i \right. \\ & \left. - 2 \left(2(1-x) - (x+2)x_W \right) x_W \right), \end{aligned} \quad (24)$$

$$f_2^2 = \frac{x}{(1-x)^2} \left(3(1-x-2(x_i+x_W))x_i + 2(6x_W-5(1-x))x_W \right), \quad (25)$$

$$f_3^2 = \frac{1}{(1-x)^2} \left((x^2+x-2-(x^2-5x-2)(x_i+x_W))x_i - 2(3x^2-x-2-(x^2-5x-2)x_W)x_W \right), \quad (26)$$

$$g_2 = \frac{1}{(1-x)^2} \left(2x((1-x)(1-x-4x_W)+3x_W^2)x_W + (3xx_i^2+(x_i-1)(1-x)^2+(2(1+2x)-3x(2x+3x_W))x_W)x_i \right), \quad (27)$$

In writing the above expressions, the unitarity condition $\sum_i V_{u_1 i} V_{t i}^\dagger = 0$ was taken into account, i.e., any term independent of the internal quark mass was dropped out. Also, it is straightforward to show that $\sum_{a=1}^3 f_a^b = 0$ for $b = 1, 2$, which means that, as expected, the A_i amplitudes are free of ultraviolet divergences.

We now are ready to calculate the contribution of the $tu_1 g^*$ vertex to the $t \rightarrow u_1 \bar{u}_2 u_2$ decay. Below, p_2 and \bar{p}_2 will stand for the 4-momenta associated with the u_2 and \bar{u}_2 quarks. It is useful to introduce the following dimensionless variable $y = (p_1 + \bar{p}_2)^2/m_t^2$. The u_1 quark mass will be retained in the phase space integral since a factor of $1/x^2$, associated with the gluon pole, enters into the $t \rightarrow u_1 \bar{u}_2 u_2$ squared amplitude. Using the expressions for the $tu_1 g^*$ and $g^* u_2 u_2$ vertices, it is straightforward to construct the invariant amplitude associated with the diagram (i) of Fig. 1. After making this, we can write the invariant mass distribution $d\Gamma/dx$ as follows

$$\frac{d\Gamma(t \rightarrow u_1 u_2 \bar{u}_2)}{dx} = \frac{m_t}{256\pi^3} \int dy \sum_{spins} |\mathcal{M}|^2, \quad (28)$$

with the squared amplitude being

$$\begin{aligned} \sum_{spins} |\mathcal{M}|^2 &= \frac{\alpha_s^2 \alpha^2}{9s_W^4 x^2} (F_1(x, y) |A_1|^2 + F_{12}(x, y) 2\text{Re}(A_1 A_2^*) \\ &\quad + F_2(x, y) |A_2|^2), \end{aligned} \quad (29)$$

As far as the $F_i(x, y)$ functions are concerned, they are given by

$$F_1(x, y) = -4(x^2 + x(2y-1) + 2y(y-1)), \quad (30)$$

$$F_{12}(x, y) = -4x(1-x), \quad (31)$$

$$F_2(x, y) = 4x(2y(y-1) + x(2y-1) + 1). \quad (32)$$

whereas the Passarino-Veltman scalar functions B_0 and C_0 are, in the usual notation:

$$B_0(1) = B_0(0, m_t^2 x_i, m_W^2), \quad (33)$$

$$B_0(2) = B_0(m_t^2 x, m_t^2 x_i, m_t^2 x_i), \quad (34)$$

$$B_0(3) = B_0(m_t^2, m_t^2 x_i, m_W^2), \quad (35)$$

$$C_0(x, x_i) = C_0(m_t^2, 0, m_t^2 x, m_t^2 x_i, m_W^2, m_t^2 x_i). \quad (36)$$

The integration limits are as follows

$$y_{min} = y_0^2 + \frac{1}{2}(1-x) \left(1 - \sqrt{1 - \frac{4y_0^2}{x}} \right), \quad (37)$$

$$y_{max} = y_0^2 + \frac{1}{2}(1-x) \left(1 + \sqrt{1 - \frac{4y_0^2}{x}} \right), \quad (38)$$

$$4y_0^2 \leq x \leq (1-x_0)^2, \quad (39)$$

where $x_0 = m_{u_1}/m_t$ and $y_0 = m_{u_2}/m_t$.

From Eq. (16) it is evident that A_1 , the monopole term, vanishes in the on-shell limit ($x \rightarrow 0$), in agreement with the fact that the $t \rightarrow u_1 g$ decay is only determined by a dipole term. We will show below that the monopole contribution is slightly larger than the dipole one.

Since the A_i amplitudes do not depend on y , this variable can be integrated over readily. In the $y_0 \rightarrow 0$ limit, we are left with

$$\frac{d\Gamma}{dx} = \frac{\alpha_s^2 \alpha^2 m_t}{1728 s_W^4 \pi^3} (f_1(x)|A_1|^2 + f_{12}(x)Re(A_1 A_2^*) + f_2(x)|A_2|^2), \quad (40)$$

where

$$f_1(x) = \frac{1}{x^2}(1+2x)(1-x)^2, \quad (41)$$

$$f_{12} = -\frac{6}{x}(1-x)^2, \quad (42)$$

$$f_2(x) = \frac{1}{x}(2+x)(1-x)^2. \quad (43)$$

It is interesting to note that we will not take into account the limiting case $y_0 \rightarrow 0$ ($m_{u_2} = 0$) when integrating over x because the $d\Gamma/dx$ distribution would become undefined in $x = 0$ due to the gluon pole. This corresponds to the case when the u_2 quark emerges parallel to \bar{u}_2 and we cannot take the limit of massless u_2 quark as it would lead to a collinear singularity. Thus, although we have neglected the outgoing quark masses in the transition amplitude, they must be retained in the integration limits of the x variable.

III. NUMERICAL RESULTS AND FINAL REMARKS

For the numerical analysis we will use the values of the running coupling constant α_s and quark masses at the m_t scale, namely, $\alpha_s(m_t) = 0.10683$, $m_t(m_t) = 174.3$ GeV, $m_b(m_t) = 2.85$ GeV, $m_c(m_t) = 0.63$ GeV, $m_s(m_t) = 0.09$ GeV, $m_d(m_t) = 0.0049$ GeV, and $m_u(m_t) = 0.00223$ GeV [15]. It is worth noting that the numerical results do not change considerably for small variations of the outgoing quark masses.

We first would like to compare the size of the off-shell $tu_1 g^*$ vertex with that of the on-shell one. Numerical evaluation shows that the $tu_1 g$ dipole contribution is two orders of magnitude smaller than the $tu_1 g^*$ monopole contribution and one order smaller than the $tu_1 g^*$ dipole contribution. Thus, while the $t \rightarrow u_1 g$ decay only receives the contribution of the dipolar term, the $t \rightarrow u_1 \bar{u}_2 u_2$ transition receives an extra contribution of the monopolar term, which is slightly larger than the dipolar contribution.

The fact that the contribution of the monopole form factor to the $tu_1 g^*$ vertex is larger than that of the dipole one is exhibited in the invariant mass distribution $d\Gamma(t \rightarrow u_1 \bar{u}_2 u_2)/dx$, which is shown in Fig. 3, where we have plotted separately the monopolar and dipolar contributions. Therefore, it is evident that the $t \rightarrow u_1 \bar{u}_2 u_2$ decay is slightly dominated by the monopolar term.

We now turn to the numerical evaluation of $Br(t \rightarrow u_1 \bar{u}_2 u_2)$. Using the values given above and $\Gamma(t \rightarrow bW) = 1.55$ GeV, we obtain

$$Br(t \rightarrow u_1 \bar{u}_2 u_2) = 3.38 \times 10^{-12}. \quad (44)$$

On the other hand, according to the literature $Br(t \rightarrow u_1 g) = 5.73 \times 10^{-12}$ [7, 16]. This result shows that $Br(t \rightarrow u_1 \bar{u}_2 u_2)$ is about of the same order of magnitude than $Br(t \rightarrow u_1 g)$. If one sums over all the possible $\bar{u}_2 u_2$ pairs, the resulting $Br(t \rightarrow u_1 \bar{u}_2 u_2)$ is of the order of 10^{-11} and thus larger than $BR(t \rightarrow u_1 g)$. Although these decay rates seem exceedingly small to be detected ever, they may be largely enhanced in some SM extensions. In such a case the effect discussed above may have some interesting implications.

In conclusion, we have shown the interesting fact that three-body decay $t \rightarrow u_1 \bar{u}_2 u_2$ has a branching ratio about the same order of magnitude than the one of the two-body decay $t \rightarrow u_1 g$. Although rare decays of this type are very suppressed in the SM, they may have much larger branching ratios in other SM extensions, thereby constituting an interesting place to search for any new physics effects.

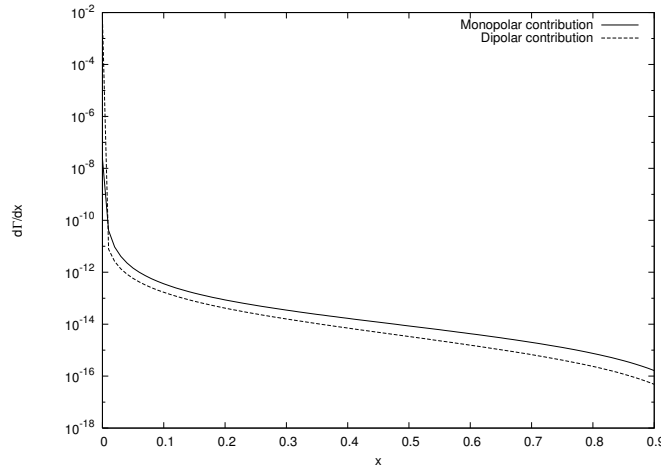


FIG. 3: Monopolar and dipolar contributions to the $t \rightarrow u_1 \bar{u}_2 u_2$ invariant mass distribution $d\Gamma/dx$.

Note added

After this work was submitted, a preprint was posted to the preprint archive by Eilam, Frank and Turan[16], who evaluate the $t \rightarrow cgg$ and $t \rightarrow u_1 \bar{u}_2 u_2$ decays. Although these authors do not present explicit analytical expressions for the $t \rightarrow u_1 \bar{u}_2 u_2$ decay [16], our numerical result for the branching ratio agrees with theirs. We have also learnt that Deshpande, Margolis and Trottier presented a similar analysis in [17]. These authors reached a similar conclusion on the $t \rightarrow c\bar{q}q$ decay in both the standard and the two-Higgs doublet models.

Acknowledgments

We acknowledge support from SNI and CONACYT under grant U44515-F.

-
- [1] F. Abe *et al.*, Phys. Rev. Lett. **74**, 2626 (1995); S. Abachi *et al.*, *ibid.* **74**, 2632 (1995).
 - [2] F. Abe *et al.*, Phys. Rev. **D51**, 4623 (1995); Phys. Rev. Lett. **80**, 2767 (1998); **80**, 2779 (1998); **82**, 271 (1999); **80**, 2773 (1998); **79**, 1992 (1997); **79**, 3585 (1997); **80**, 5720 (1998); **80**, 2525 (1998); S. Abachi *et al.*, *ibid.* **79**, 1197 (1997); **79**, 1203 (1997); B. Abbott *et al.*, *ibid.* **80**, 2063 (1998); **83**, 1908 (1999); Phys. Rev. **D58**, 052001 (1999); **D60**, 052001 (1999); T. Affolder *et al.*, Phys. Rev. Lett. **84**, 216 (2000).
 - [3] For a recent review on top quark physics, see D. Chakraborty, J. Konigsberg, D. Rainwater, Ann. Rev. Nucl. Part. Sci. **53**, 301 (2003). See also M. Beneke *et al.*, arXiv:hep-ph/0003033.
 - [4] Elizabeth Jenkins, Phys. Rev. **D56**, 458 (1997).
 - [5] R. Decker, M. Nowakowski and A. Pilaftsis, Z. Phys. C **57**, 339 (1993); G. Mahlon and S. J. Parke, Phys. Lett. B **347**, 394 (1995).
 - [6] S. Eidelman *et al.*, Phys. Lett. **B592**, 1 (2004).
 - [7] G. Eilam, J. L. Hewett, A. Soni, *ibid.* **D44**, 1473 (1991); 59, 039901(E) (1999).
 - [8] J. L. Díaz-Cruz, R. Martínez, M. A. Pérez, and A. Rosado, Phys. Rev. **D41**, 891 (1999).
 - [9] B. Mele and S. Petrarca, Phys. Lett. **B435**, 401 (1999).
 - [10] M. E. Luke, M. J. Savage, Phys. Lett. **B307**, 387 (1993); D. Atwood, L. Reina, A. Soni, Phys. Rev. **D53**, 1199 (1996); Phys. Rev. Lett. **75**, 3800 (1995); E. O. Iltan, Phys. Rev. **D65**, 075017 (2002); E. O. Iltan and I. Turan, Phys. Rev. **D67**, 015004 (2003); W. S. Hou, Phys. Lett. **296**, 179 (1992); J. L. Díaz-Cruz, M. A. Pérez, G. Tavares-Velasco, J. J. Toscano, Phys. Rev. **D60**, 115014 (1999); D. Atwood, L. Reina, A. Soni, Phys. Rev. **D55**, 3156 (1997).
 - [11] G. M. de Divitiis, R. Petronzio, L. Silvestrini, Nucl. Phys. **B504**, 45 (1997); J. L. Lopez, D. V. Nanopoulos, R. Rangarajan, Phys. Rev. **D56**, 3100 (1997); C. S. Li, R. J. Oakes, J. M. Yang, Phys. Rev. **D49**, 293 (1994); J. Yang, C. S. Li, Phys. Rev. **D49**, 3412 (1994); G. Couture, C. Hamzaoui, H. Konig, Phys. Rev. **D52**, 1713 (1995); G. Couture, M. Frank, H. Konig, Phys. Rev. **D56**, 4213 (1997); Jun-jie Cao, Zhao-hua Xiong, Jin Min Yang, Nucl. Phys. **B651**, 87 (2003); Jian Jun Liu, Chong Sheng Li, Li Lin Yang, Li Gang Jin, Phys. Lett. B **599**, 92 (2004).
 - [12] J. M. Yang, C. S. Li, Phys. Rev. **D49**, 3412 (1994); J. M. Yang, B. -L. Young, X. Zhang, Phys. Rev. **D58**, 055001 (1998); G. Eilam, A. Gemintern, T. Han, J. M. Yang, X. Zhang, Phys. Lett. **B510**, 227 (2001).

- [13] Chong-xing Yue, Gong-ru Lu, Qing-jun Xu, Guo-li Liu, Guang-ping Gao, Phys. Lett. **B508**, 290 (2001); J. A. Aguilar-Saavedra, B. M. Nobre, Phys. Lett. **B553**, 251 (2003); Gong-ru Lu, Fu-rong Yin, Xue-lei Wang, Ling-de Wan, Phys. Rev. **D68**, 015002 (2003); R. Gaitan, O. G. Miranda and L. G. Cabral-Rosetti, Phys. Rev. D **72**, 034018 (2005).
- [14] F. del Aguila, M. Pérez-Victoria, J. Santiago, Phys. Lett. **B492**, 98 (2000); J. High Energy Phys. 09, 011 (2000); Adriana Cordero-Cid, M. A. Pérez, G. Tavares-Velasco, J. J. Toscano, Phys. Rev. **D70**, 074003 (2004).
- [15] H. Fusaoka and Y. Koide, Phys. Rev. D **57**, 3986 (1998).
- [16] G. Eilam, M. Frank, and I. Turan, arXiv:hep-ph/0601151.
- [17] N. G. Deshpande, B. Margolis, and H. D. Trottier, Phys. Rev. D **45**, 178 (1992).